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New developments in inflationary models[1]

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Abstract

We study models of inflation where the inflaton corresponds to a flat direction in field space and its mass term is generated by gravity mediated soft supersymmetry breaking at high scale. Assuming the inflaton to have non negligible couplings to other fields, its mass runs with scale and can reach the small value required by slow roll inflation at a lower scale, even if its initial value is too large. Slow roll inflation can therefore take place in such a regime, as long as the mass remains small, with a spectral index that is then scale dependent. We explore the parameter space of this kind of models to find the region compatible with the present observations.

1 Introduction

Slow roll inflation[2] requires the flatness conditions $\epsilon \ll 1$ and $|\eta| \ll 1$ on the potential, where

$$\epsilon \equiv \frac{1}{2} M_{Pl}^2 \left(\frac{V'}{V}\right)^2 ; \;\; \eta \equiv M_{Pl}^2 \frac{V''}{V} , \tag{1}$$

and $M_{Pl} \equiv (8\pi G)^{-1/2} = 2.4 \times 10^{18}$GeV. The first condition is easily satisfied in most inflationary models where the inflaton lies near to an extremum of the potential, but the second is problematic. In fact, during inflation in a generic supergravity potential all scalar fields [2, 3], and in particular the inflaton [4] acquire a contribution to the mass-squared of magnitude V/M_{Pl}^2, which spoils this condition.

Very few proposals have been put forward to solve this problem which would not be relying on some sort of fine tuning [1]; specific types of Kähler potential and superpotential can succeed in canceling the dangerous contribution, as it

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[2] For a general discussion and references on inflation see [1]

happens in no–scale supergravity. The aim of this paper is to investigate instead the proposal of Stewart [5, 6]: in this scenario the contribution to the inflaton mass is unsuppressed at high scale, but loop corrections can flatten the inflaton potential to realize sufficient inflation without any significant fine-tuning.

We will explore the potential in a general model of this kind and find the region of parameter space allowed by the observed magnitude and spectral index of the curvature perturbation. We will finally discuss the naturalness of such picture and the consequence of future observations.

2 The running mass models

In the model proposed by Stewart [5], slow-roll inflation occurs, with the following Renormalization Group (RG) improved potential for the canonically normalized inflaton field ϕ;

$$V = V_0 + \frac{1}{2} m_\phi^2(\phi)\phi^2 + \frac{1}{2} m_\psi^2(\phi)\psi^2 + \frac{1}{4}\lambda(\phi)\phi^2\psi^2 \cdots . \tag{2}$$

The constant term V_0 is supposed to come from the supersymmetry breaking and to dominate at all relevant field values. Non–renormalizable terms, represented by the dots, give the potential a minimum at large ϕ, but they are supposed to be negligible during inflation. The last two terms also vanish during inflation, since $\psi = 0$, but are responsible for the hybrid exit from the inflationary period.

The inflaton mass-squared and all the other parameters depend on the renormalization scale Q, and following [5, 6] we have taken $Q = \phi$, where now ϕ denotes the classical v.e.v of the inflaton field during inflation. Such choice for the renormalization scale minimizes the one loop correction to the potential, since the main contribution goes like $\ln(\phi/Q)$ for ϕ larger than any other scale, and therefore the potential in eq. (2) is effectively equivalent to the full one loop potential. If the inflaton v.e.v. is not the dominant scale, then some other choice of Q will be appropriate and the simplification we have made is no more viable. We will assume that the inflaton v.e.v. is the dominant scale up to the end of inflation.

At the Planck scale, $m_\phi^2(M_{Pl})$ is supposed to have the generic magnitude

$$|m_0^2| = |m_\phi^2(M_{Pl})| \sim V_0 \tag{3}$$

coming from supergravity corrections [4, 7].

Without running, this would give $|\eta| \sim 1$, preventing slow–roll inflation. But at field values below the Planck scale, the RG drives $m_\phi^2(\phi)$ to small values, corresponding to $|\eta(\phi)| \ll 1$, and slow–roll inflation can take place. We have in fact that the slow-roll parameters are given in our case by

$$\epsilon = \frac{M_{Pl}^2\phi^2}{2V_0^2}\left[m_\phi^2(\phi) + \frac{1}{2}\frac{dm_\phi^2}{d\ln(\phi)}\right]^2 \tag{4}$$

$$\eta = \frac{M_{Pl}^2}{V_0}\left[m_\phi^2(\phi) + \frac{3}{2}\frac{dm_\phi^2}{d\ln(\phi)} + \frac{1}{2}\frac{d^2 m_\phi^2}{d\ln^2(\phi)}\right]; \tag{5}$$

since the derivatives of m_ϕ^2 are suppressed by the coupling constant, both ϵ and η are small around the value of ϕ where the inflaton mass vanishes and inflation can successfully happen in such conditions.

Since in this model the η parameter changes considerably as ϕ decreases, slow-roll inflation will continue until some epoch ϕ_{end}, when either the critical value $\phi_c = -2m_\psi/\lambda$ is reached or $\eta(\phi)$ becomes of order 1. To reduce the number of parameters involved in our analysis, we will assume the latter to be the case, so that both the flattening of the potential and the end of slow roll inflation are due to the mass running; the critical value ϕ_c will be reached after a brief phase of fast–roll, that should not change considerably the e-foldings number. We have then that the number of e-folds generated while the inflaton runs from value ϕ to ϕ_{end} is given in the slow roll approximation by

$$\begin{aligned} \mathcal{N}(\phi) &= \int_{\phi_{end}}^{\phi} d\phi \frac{V}{M_{Pl}^2 V'} \\ &= \frac{V_0}{M_{Pl}^2}\int_{\phi_{end}}^{\phi} \frac{d\ln(\phi)}{m_\phi^2(\phi) + \frac{1}{2}\frac{dm_\phi^2(\phi)}{d\ln(\phi)}} \end{aligned} \tag{6}$$

where $m_\phi^2(\phi)$ is given by solving the RG equations.

3 The scale dependence of the spectral index

The scale dependence of the spectral index in this kind of models is strictly related to the RG equation of the inflaton mass. The case of a gauge coupling dominated running has been studied in [6, 7, 8] while the Yukawa dominated running has been considered in [9]. We will review the two extreme cases in a toy model in the following, concentrating in the case where inflation takes place while the inflaton rolls from the region $m_\phi^2 \simeq 0$ towards the origin.

A useful way to understand the general behaviour, is to consider the linear approximation for the running inflaton mass

$$m_\phi^2(\phi) \simeq -\frac{V_0}{M_{Pl}^2}\left[\mu_\star^2 + c\ln(\phi/\phi_\star)\right] \tag{7}$$

where the $\star$ denotes values of the variables where V' vanishes and the constant c is small and proportional to the relevant coupling. We have then that the observational quantities can all be written as function of the three adimensional parameters

$$c = -\frac{dm_\phi^2}{d\ln(\phi)}|_{\phi=\phi_\star} = -2\mu_\star^2 \tag{8}$$

$$\begin{aligned}\tau &= -|c|\ln(\phi_\star/M_{Pl}) &(9)\\ \sigma &= \lim_{\phi\to\phi_\star} ce^{c\mathcal{N}(\phi)}\ln(\phi_\star/\phi), &(10)\end{aligned}$$

where $\mathcal{N}(\phi)$ is given by eq. (6). Assuming the linear approximation to hold up to the end of inflation, such expression simplifies to

$$\begin{aligned}\mathcal{N}(\phi) &= -\frac{1}{c}\int_{\phi_{end}}^{\phi}\frac{d\ln(\phi)}{\ln(\phi/\phi_\star)} &(11)\\ &= -\frac{1}{c}\ln\left|\frac{\ln(\phi/\phi_\star)}{\ln(\phi_{end}/\phi_\star)}\right|. &(12)\end{aligned}$$

Note that the e-folding number is inversely proportional to the coupling (contained in c) so that a small coupling gives automatically sufficient inflation. We see also that the parameter σ gives directly a measure of the departure from the linear approximation at ϕ_{end}, since in the case this holds all the way, σ should be given by $\sigma = \pm 1 + c$ for $\phi \to \phi_\star$, i.e. would be a number of order 1. In general such approximation breaks down well before ϕ_{end} and σ can take also very large values.

We obtain then for the spectral index in the linear approximation:

$$n(\mathcal{N}) - 1 = 2\sigma e^{-c\mathcal{N}} - 2c; \tag{13}$$

while the COBE normalization imposes a constraint on V_0:

$$\frac{V_0^{1/2}}{M_{Pl}^2} = 5.3\times 10^{-4}|\sigma|\exp\left[-\frac{\tau}{c} - c\mathcal{N}_{COBE} - \frac{\sigma}{c}e^{-c\mathcal{N}_{COBE}}\right]. \tag{14}$$

So for every particular model, the experimental constraints limit the range of the parameters allowed. We see in Fig. 1 the region in the $\sigma - c$ plane compatible with a spectral index $|n-1| \leq 0.2$ at $\mathcal{N} = 50$.

Note that every quadrant correspond to a different type of inflationary model:

- positive c implies that the inflaton mass changes sign from negative to positive and therefore a maximum in the RG improved potential develops around $m_\phi^2 \simeq 0$; in such case the sign of sigma indicates if the inflaton is rolling towards the origin ($\sigma > 0$) or towards large fields values ($\sigma < 0$), always away from the maximum;

- negative c implies on the contrary that the inflaton mass changes sign from positive to negative and a minimum in the RG improved potential develops around $m_\phi^2 \simeq 0$; again the sign of sigma is related to the direction of the inflaton's motion: $\sigma > 0$ means that the inflaton is rolling towards the minimum from the large field values while $\sigma < 0$ from the small field values.

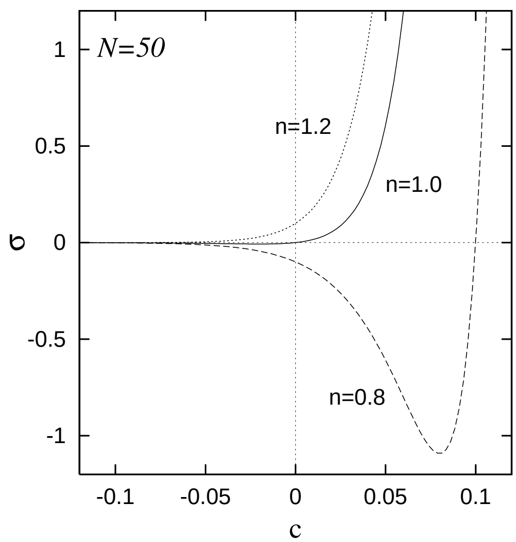

Figure 1: Lines of constant spectral index in the $\sigma - c$ plane for $\mathcal{N} = 50$; assuming that this e-folding number corresponds to COBE scales the allowed region is between the long–dashed and dotted lines.

The COBE normalization gives additional bounds on the value of τ for every choice of c, σ, since V_0 has to be surely larger that the scale of nucleosynthesis (for instant reheating we have in fact $T_{RH} = V_0^{1/4}$, but usually $T_{RH} < V_0^{1/4}$).

4 A simple toy model

Let us consider as an example the case of the superpotential

$$W = \lambda S \mathrm{Tr}\,(\phi_1 \phi_2) \tag{15}$$

where S is a singlet chiral superfield, while ϕ_i are chiral superfield in the adjoint representation of the gauge group $SU(N)$. We can easily compute the scalar potential given by (15) in the limit of unbroken supersymmetry and, writing the adjoint fields in the fundamental basis[3] $\phi_i = \phi_i^a t_a$, it is given by:

$$V = \frac{\lambda^2}{4} |\phi_1^a \phi_2^a|^2 + \frac{\lambda^2}{4} |S|^2 (|\phi_1^a|^2 + |\phi_2^a|^2) + \frac{|D_a|^2}{2} \tag{16}$$

[3] We define the fundamental representation of $SU(N)$ t_a such that $\mathrm{Tr}\,(t_a t_b) = \frac{1}{2}\delta_{ab}$ and $[t_a, t_b] = f_{abc} T_c$, while for the adjoint representation, e.g. $T^a_{ij} = f_{aij}$, we have $\mathrm{Tr}\,(T_a T_b) = N\delta_{ab}$.

where S, ϕ_i indicate now the scalar components of the chiral multiplets, summation over a is implicit and

$$D_a = i\frac{g}{2} f_{abc} \left(\phi_1^{b*} \phi_1^c + \phi_2^{b*} \phi_2^c \right) \tag{17}$$

with g denoting the $SU(N)$ gauge coupling.

We see clearly that a flat direction exists for

$$\begin{aligned} S &= 0 \qquad (18)\\ \phi_1^a \phi_2^a &= 0 \qquad (19)\\ f_{abc}\, \phi_i^{b*} \phi_i^c &= 0. \qquad (20) \end{aligned}$$

This is not the most general case and other flat directions are present, parameterized by gauge invariant polynomials [12].

In the following we will consider the case when the inflaton is one of the components of the charged fields, i.e. we will take $\phi = Re\,[\phi_1^a]$ to be the inflaton, while all the other fields are supposed to vanish during inflation.

Then the potential for the inflaton is reduced only to the soft susy breaking terms [10] and assumes the form of eq. (2), where V_0 is a cosmological constant that is generated by some other sector of the theory and is canceled in the true vacuum by the v.e.v. of a field in our sector, playing the role of the ψ.

From supergravity, we expect all the susy breaking scalar masses, respectively $m_S, m_{1/2}$ for the singlet and charged fields, to be of order of $V_0^{1/2}/M_{Pl}$ and the trilinear parameter Y, in $\frac{Y}{2}\lambda S \phi_1^a \phi_2^a + h.c.$, to be of the same order, as $V_0^{1/4}$ is the scale of explicit supersymmetry breaking during inflation. Note however that while the contribution to the scalar masses coming from V_0 is always present, the trilinear coupling Y_0 not always receives a contribution proportional to $V_0^{1/2}$. Moreover, at the end of inflation V_0 vanishes and the susy breaking parameters will be connected instead to the gravitino mass in the usual way, so in principle the susy breaking parameters during and after inflation are different.

In order to write the RG improved potential, we will need to consider the one loop renormalization group equations for all our parameters and extract the behaviour of the inflaton mass.

Following [11], we write down the equations for our particle content. The gauge field strength $\alpha = g^2/(4\pi)$ and the gaugino mass satisfy

$$\begin{aligned} \frac{d\alpha}{dt} &= \frac{\beta}{2\pi}\alpha^2 \qquad (21)\\ \frac{d\tilde{m}}{dt} &= \frac{\beta}{2\pi}\alpha\tilde{m} \qquad (22) \end{aligned}$$

where $t = \ln(Q)$ is the renormalization scale and $\beta = -N$ in our case of $SU(N)$ with two matter superfields in the adjoint representation ($\beta = -3N + n_{adj}N$).

This two equations are independent from the others and their solution is

$$\alpha(t) = \frac{\alpha_0}{1 - \frac{\beta}{2\pi}\alpha_0 t} = \frac{\alpha_0}{1 + \tilde{\alpha}_0 t} \tag{23}$$

$$\tilde{m}(t) = \frac{\tilde{m}_0}{\alpha_0}\alpha(t) \tag{24}$$

where $\tilde{\alpha}_0 = N\alpha_0/(2\pi)$ and a 0 subscript denotes quantities at the Planck scale.

For the Yukawa coupling, which we can always take real absorbing its phase in the definition of the singlet field S, we have instead

$$\frac{d\lambda}{dt} = -N\frac{\alpha}{\pi}\lambda + \frac{\lambda}{16\pi^2}(N^2+1)|\lambda|^2 \tag{25}$$

while for the soft susy breaking masses the equations can be cast in a simple form using the variables

$$m^2_{1-2} = m^2_1 - m^2_2, \tag{26}$$

$$m^2_{1-S} = m^2_1 - \frac{1}{N^2-1}m^2_S \tag{27}$$

and m^2_S, where m_S, m_i are respectively the susy breaking masses of S, ϕ_i and Y is the susy breaking trilinear coupling. In fact we have:

$$\frac{dm^2_{1-2}}{dt} = 0 \tag{28}$$

$$\frac{dm^2_{1-S}}{dt} = -\frac{2N\alpha}{\pi}\tilde{m}^2 \tag{29}$$

$$\frac{dm^2_S}{dt} = \frac{N^2+1}{8\pi^2}|\lambda|^2 m^2_S + \frac{N^2-1}{8\pi^2}|\lambda|^2\left[2m^2_{1-S} - m^2_{1-2} + \frac{|Y|^2}{2}\right]. \tag{30}$$

The trilinear term will have instead the equation

$$\frac{dY}{dt} = \frac{1}{32\pi^2}(N^2+1)Y|\lambda|^2 + \frac{2}{\pi}N\alpha\tilde{m}. \tag{31}$$

These are a system of coupled differential equations. We will in the next sections consider approximate solutions in different cases and obtain the running inflaton mass.

5 Dominant gauge coupling

For $\alpha \gg \lambda^2$ a model independent analysis has been made in [7]. In this case we can neglect the λ^2 terms and the inflaton mass running does not depend on the Yukawa coupling.

We have then for both charged fields

$$m_i^2(t) = m_{i,0}^2 - 2\tilde{m}_0^2 \left[1 - \frac{1}{(1+\tilde{\alpha}_0 t)^2}\right]. \tag{32}$$

Notice that eq.(32) gives in general a solution of eq.(29), in the case of non negligible Yukawa.

In this case we can easily translate our three parameters into physical ones and we have:

$$\begin{aligned}
c &= 2\tilde{\alpha}_0 A_0 \left[1 + \frac{\mu_0^2}{A_0}\right]^{3/2} && (33)\\
\tau &= 2A_0 \left(1 + \frac{\mu_0^2}{A_0}\right)\left[\sqrt{1 + \frac{\mu_0^2}{A_0}} - 1\right] && (34)\\
\ln(\sigma) &= 2\left(1 + \frac{\mu_0^2}{A_0}\right)\left[\frac{1}{\sqrt{1+\frac{\mu_0^2}{A_0}}} - \frac{1}{\sqrt{1+\frac{\mu_0^2+1}{A_0}}}\right] \\
&\quad + \ln\left[4\left(A_0 + \mu_0^2\right)\right] + \ln\left[\frac{\sqrt{1+\frac{1}{\mu_0^2+A_0}} - 1}{\sqrt{1+\frac{1}{\mu_0^2+A_0}} + 1}\right] && (35)
\end{aligned}$$

where $\mu_0^2 = |m_{1,0}^2| M_{Pl}^2/V_0$ and $A_0 = 2\tilde{m}_0^2 M_{Pl}^2/V_0$ are the values of the scalar and gaugino masses at the Planck scale. Then the allowed region for the c, σ parameters shown in the first quadrant of Fig. 1 gives the bounds on physical parameters shown in Fig. 2.

Notice that in this particular case, an inflaton mass squared of order V_0/M_{Pl}^2 at the Planck scale is acceptable, and the running is efficiently flattening the potential, provided that the gaugino mass is sufficiently large. Notice that, as shown in the graph for a specific choice of the gauge coupling, generally gaugino masses larger than the scalar one have to be assumed.

For consistency we have also to find the range of values of λ where this approximation is reliable: naturally the limit $\lambda \to 0$ violates our assumption $\phi_{end} > \phi_c$, so that we have a lower bound on λ:

$$\lambda_0^2 \geq 4V_0 \exp\left[\frac{2}{\tilde{\alpha}_0}\left(1 - \frac{1}{\sqrt{1+\frac{V_0+|m_{1,0}^2|}{A_0}}}\right)\right]. \tag{36}$$

As we can see this bound is very sensitive to the value of the gauge coupling and also V_0; in general α_0 has to be of the order of 0.01 or so to give a non negligible allowed region for the initial masses and a non negligible allowed range for λ.

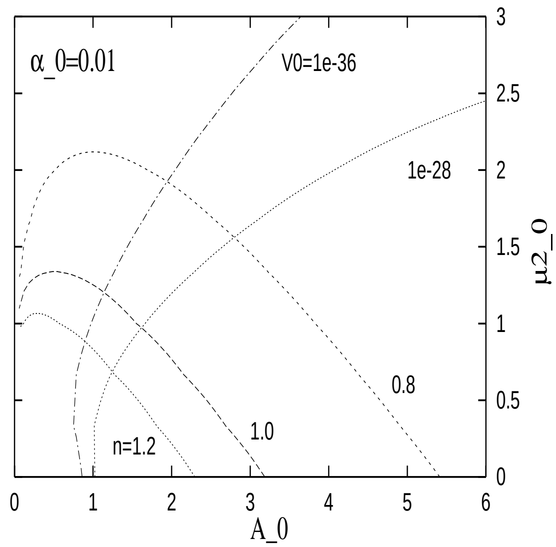

Figure 2: Lines of constant spectral index in the case of gauge dominated running of the inflaton mass in the plane $\mu_0^2 = |m_\phi^2(M_{Pl})|M_{Pl}^2/V_0$ vs $A_0 = 2\tilde{m}^2(M_{Pl})M_{Pl}^2/V_0$ for a gauge coupling $\tilde{\alpha}_0 = N/(2\pi)\,\alpha(M_{Pl}) = 0.01$. Also the lines of constant V_0 are displayed in units of M_{Pl}^4 for $\mathcal{N}_{COBE} = 45$. This region corresponds to positive σ and c [8] .

6 Dominant Yukawa coupling

In this case the equations become similar to those for uncharged fields. We can therefore consider at the same time the model where ϕ_i are just two singlet fields substituting in the following $N^2 \to 2$. This substitution amounts to consider only one degree of freedom instead of the $N^2 - 1$ of a field in the adjoint representation of $SU(N)$. The solutions for the scalar masses are given by:

$$\begin{aligned} m_S^2(t) &= \frac{N^2-1}{N^2+1}\Bigg[(m_{S,0}^2 + m_{1,0}^2 + m_{2,0}^2 + Y_0^2)\frac{1}{1-\tilde{\lambda}_0^2 t} \\ &\quad -Y_0^2\frac{1}{\sqrt{1-\tilde{\lambda}_0^2 t}} - m_{1,0}^2 - m_{2,0}^2 + \frac{2}{N^2-1}m_{S,0}^2\Bigg] \qquad (37) \\ m_i^2(t) &= m_{i,0}^2 + \frac{1}{N^2-1}(m_S^2(t) - m_{S,0}^2), \qquad (38) \end{aligned}$$

where the subscript 0 again indicates the initial values (defined at the Planck scale) and $\tilde{\lambda}_0^2 = \frac{N^2+1}{2\pi^2}\lambda_0^2$.

We can see that in this case the initial conditions determine whether one of the masses changes sign and an extremum in the potential is reached; generally the singlet mass appears to have the stronger running since it interacts with more fields. One possibility for inflation in this case, is to have negative initial masses and chose the inflaton to be the field whose mass first become positive. In the case of universal initial masses such a field turns out to be the singlet; we have then that inflation happens in the S direction[4] and we can easily compute the parameters c, σ, τ for vanishing Y_0:

$$c = \frac{4(N^2-2)^2}{3N^4-1}\tilde{\lambda}_0\mu_0^2 \tag{39}$$

$$\tau = \frac{2(N^2-2)}{3(N^2-1)}\mu_0^2 \tag{40}$$

$$\ln(\sigma) = \frac{1}{\frac{2(N^2-2)}{N^2+1}\mu_0^2 - 1} - \ln\left[1 - \frac{N^2+1}{2(N^2-2)\mu_0^2}\right], \tag{41}$$

where $\mu_0^2 = |m_S^2(M_{Pl})|M_{Pl}^2/V_0$. In this case we have only two physical parameters to play with, but an acceptable region exists as shown in Fig. 3: for initial values of μ_0^2 of order 1, a Yukawa coupling of order 0.05 is needed to flatten the potential and V_0 is fixed by the COBE normalization to be of order $10^{12-14} GeV$.

Notice that, in contrast with the previous case of dominant gauge coupling, now μ_0^2 plays both the role of the scalar mass and of the gaugino mass and therefore a large initial μ_0^2 is needed in order for the running to be efficient. As plotted in Fig. 3, μ_0^2 has to be larger than 0.5, otherwise the η parameter never becomes of order 1 and the end of inflation has to be defined by the critical value s_c. In such a case the expression for σ is much more involved and we will not consider it.

Inflation is also possible along the charged field direction, but only for non universal masses and low values of N. Taking as an example the case $N = 2$, we have in terms of the physical quantities, $\mu_0^2 = |m_{i,0}^2|M_{Pl}^2/V_0$ and $\xi_0^2 = m_{i,0}^2/m_{S,0}^2$:

$$c = \frac{(\xi_0^2-3)^2}{5(\xi_0^2+2)}\tilde{\lambda}_0\mu_0^2 \tag{42}$$

$$\tau = \frac{\xi_0^2-3}{\xi_0^2+2}\mu_0^2 \tag{43}$$

$$\ln(\sigma) = \frac{1}{\frac{\xi_0^2-3}{5}\mu_0^2 - 1} - \ln\left[1 - \frac{5}{(\xi_0^2-3)\mu_0^2}\right]. \tag{44}$$

In such case $\xi_0^2 > 3$, i.e. a singlet mass larger than the charged fields mass, is needed for flattening the potential (resembling the gaugino mass larger than

[4] Note that also the direction $S \neq 0, \phi_i^a = 0$ is a flat direction of our potential.

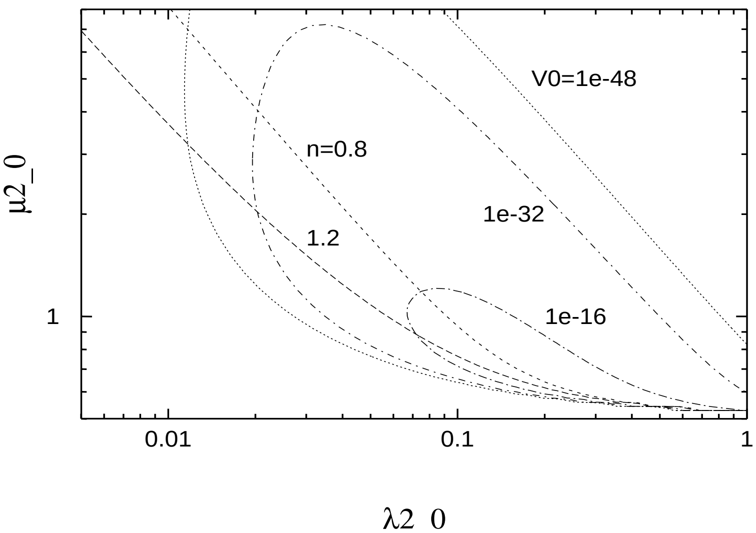

Figure 3: Lines of constant spectral index in the case of Yukawa dominated running of the inflaton mass in the plane $\mu_0^2 = |m_S^2(M_{Pl})| M_{Pl}^2/V_0$ vs $\tilde{\lambda}_0^2$ for $N \gg 1$ and $\mathcal{N}_{COBE} = 45$. Also the lines of constant V_0 are displayed in units of M_{Pl}^4. This region again corresponds to positive σ and c.

scalar mass requirement for the gauge dominated case). The bounds on the physical parameter for this non universal case are given in Fig. 4.

Another option is that of universal initial positive masses driven negative, or very small for what regards the inflaton, by the Yukawa coupling like in the case of the radiative EW breaking in the MSSM. In such a picture not only would the quantum corrections be responsible for the flattening of the potential, but also for the triggering of the hybrid–type end of inflation. Such a scenario would correspond to the quadrants with negative c in Fig. 1.

7 Conclusions

Quantum corrections can be strong enough to cancel the supergravity contribution of order V_0/M_{Pl}^2 to the inflaton mass and allow slow roll inflation to take place. The requirement to have the spectral index in the experimental range constraints tightly the parameter space of the specific models. Surprisingly anyway viable regions of the parameter space exists for reasonable values of the couplings in the different scenarios of gauge coupling or Yukawa coupling

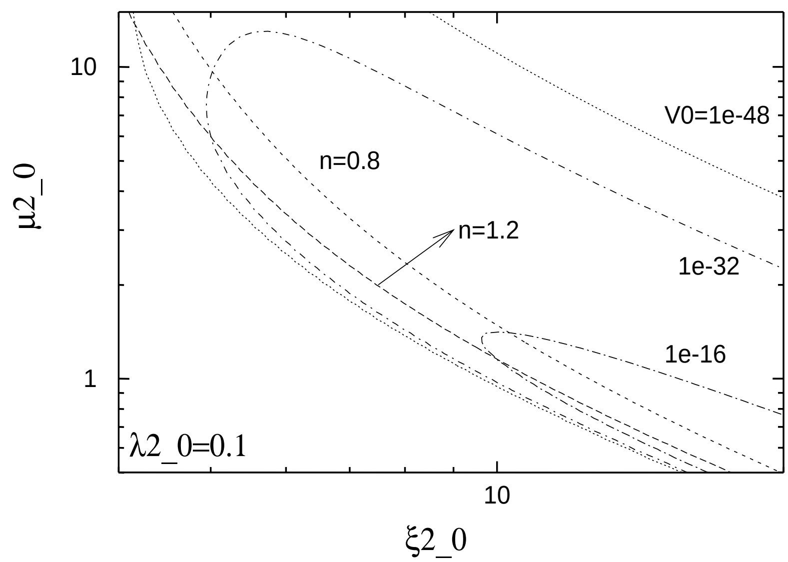

Figure 4: Lines of constant spectral index in the case of Yukawa dominated running of the inflaton mass for the non universal case. It is shown the plane $\mu_0^2 = |m_{i,0}^2| M_{Pl}^2 / V_0$ vs $\xi_0^2 = m_{S,0}^2 / m_{i,0}^2$ for $N = 2$, $\tilde{\lambda}_0^2 = 0.1$ and $\mathcal{N}_{COBE} = 45$. Also the lines of constant V_0 are displayed in units of M_{Pl}^4.

dominance.

Since the running of the inflaton mass has to be substantial to give the cancelation, in this class of models the spectral index has a significant variation on cosmological scales, surely within the reach of the Planck satellite [13]. For example in the case of gauge coupling dominance the spectral index changes by 0.1 or so in the ten e-foldings corresponding to cosmological scales [7] and such a large variation could be observed or excluded even before the launch of Planck, by the improvement of the data on the power spectrum of density perturbations. The scale dependence of the spectral index can be parameterized by eq. (13), assuming the linear approximation to be valid when cosmological scales left the horizon.

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